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ON INTRINSIC MAGNETIC MOMENTS IN BLACK HOLE CANDIDATES

Stanley L. Robertson[1] and Darryl J. Leiter[2]

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ABSTRACT

In previous work we found that many of the spectral properties of low mass x-ray binaries, including galactic black hole candidates could be explained by a magnetic propeller model that requires an intrinsically magnetized central object. Here we describe how the Einstein field equations of General Relativity and equipartition magnetic fields permit the existence of highly red shifted, extremely long lived, collapsing, radiating objects. We examine the properties of these collapsed objects and discuss characteristics that might lead to their confirmation as the source of black hole candidate phenomena.

Subject headings: black hole physics–magnetic fields–X-rays: binaries

1. INTRODUCTION

In earlier work (Robertson & Leiter 2002) we extended analyses of magnetic propeller effects (Campana et al. 1998, Zhang, Yu & Zhang 1998) of neutron stars (NS) in low mass x-ray binaries (LMXB) to the domain of galactic black hole candidates (GBHC). From the luminosities at the low/high spectral state transitions, accurate rates of spin were found for NS and accurate quiescent luminosities were calculated for both NS and GBHC. NS magnetic moments were in agreement with those found for similarly spinning 200 - 600 Hz pulsars. GBHC spins were found to be typically 10 - 50 Hz. Their magnetic moments of $\sim 10^{29}$ gauss cm^3 are ~ 100 times larger than those of 'atoll' class NS. In the magnetic propeller model, the inner disk radius, r, determines the spectral state. Very low to quiescent states correspond to an inner accretion disk radius outside the light cylinder. The inner disk radius lies between light cylinder and co-rotation radius in the low/hard/radio-loud/jet-producing state of the active propeller regime. The high/soft state corresponds to an inner disk inside the co-rotation radius and accreting matter impinging on the central object. We show here that this permits a quantitative accounting for the 'ultrasoft' high state spectral peak and a high state hard x-ray spectral tail.

A field in excess of 10^8 G has been found at the base of the jets of GRS 1915+105 (Gliozzi, Bodo & Ghisellini 1999, Vadawale, Rao & Chakrabarti 2001). A recent study of optical polarization of Cygnus X-1 in its low state (Gnedin et al. 2003) has found a slow GBHC spin and a magnetic field of $\sim 10^8$ gauss at the location of its optical emission. Given the r^{-3} dependence of field strength on magnetic moment, the implied magnetic moments are in good agreement with those we have found. Although Gnedin et al. attempted to explain the Cygnus X-1 magnetic field as a result of a spinning charged black hole, the necessary charge of 5×10^{28} esu would not be stable. Given the charge/mass ratios of electrons and protons, the opposing electric forces on them would then be at least 10^6 times the gravitational attraction of $\sim 10M_\odot$. Due to highly variable accretion rates, it is also unlikely that disk dynamos could produce the stability of fields needed to account for either spectral state switches or quiescent spin-down luminosities. Both also require magnetic fields co-rotating with the central object.

Considering the magnetic moments to be intrinsic to the central object permits a physically obvious and unified explanation of LMXB radio and spectral states, but this is incompatible with the event horizons of black hole models of the GBHC. The success of the magnetic propeller model for GBHC and the lack of evidence for event horizons in GBHC (Abramowicz, Kluzniak & Lasota 2002) strongly suggests that it must be possible, within the confines of Einstein's General Relativity to accommodate intrinsic magnetic moments in gravitationally collapsed objects. This can be achieved if the energy momentum tensor on the right hand side of the Einstein equation

$$G^{\mu\nu} = (8\pi G/c^4)T^{\mu\nu} \tag{1}$$

is chosen in a manner that dynamically enforces the Strong Principle of Equivalence (SPOE) requirement of 'timelike worldline completeness'; i.e., the requirement that the worldlines of physical matter, under the influence of both gravitational and non-gravitational forces, must remain timelike in all of spacetime (Wheeler & Ciuofolini 1995). When this SPOE condition is met, trapped surfaces leading to event horizons cannot be dynamically formed and intrinsic magnetic moments can exist in gravitationally collapsing objects (Leiter & Robertson 2003, Mitra 2000, 2002, see below).

2. MAGNETOSPHERIC, ETERNALLY COLLAPSING OBJECTS (MECO)

A relatively simple example of a collapsing, compact object that can dynamically obey the SPOE requirement of 'timelike worldline completeness' is that of a radiating plasma containing an equipartition magnetic dipole field that drives it to radiate at its Eddington limit. Such an object can be described to first order by the energy-momentum tensor:

$$T^\nu_\mu = (\rho + P/c^2)u_\mu u^\nu - P\delta^\nu_\mu + E^\nu_\mu \tag{2}$$

where $E^\nu_\mu = qk_\mu k^\nu$, $k_\mu k^\mu = 0$ describes outgoing radiation in a geometric optics approximation, ρ is energy density of matter, P is the pressure and q the flux of

[1] Physics Dept.,Southwestern Oklahoma State University, Weatherford, OK 73096 (roberts@swosu.edu)
[2] FSTC, Charlottesville, VA 22901 (dleiter@aol.com)

photon radiation. For the collapsing mass, we use a comoving interior metric given by

$$ds^2 = A(r,t)^2c^2dt^2 - B(r,t)^2dr^2 - R(r,t)^2(d\theta^2 + sin^2\theta d\phi^2) \quad (3)$$

and a non-singular exterior Vaidya metric with outgoing radiation

$$ds^2 = (1 - 2GM/c^2R)c^2du^2 + 2cdudR - R^2(d\theta^2 + sin^2\theta d\phi^2) \quad (4)$$

where R is the areal radius and $u = t - R/c$ is the retarded observer time. In order to maintain timelike worldline completeness as required by the SPOE, the surface redshift must remain finite (Leiter & Robertson 2003, Mitra 2000, 2002). Then the proper time $d\tau_s$, at the collapsing, radiating surface, S, will be positive definite if

$$d\tau_s = \frac{du}{1+z_s} = du(\Gamma_s + U_s/c) \ > \ 0 \quad (5)$$

where $U_s = dR/d\tau$ is the proper time rate of change of $R(r,t)$ and

$$\Gamma_s = (1 - \frac{2GM(r,t)_s}{c^2R_s} + \frac{U_s}{c})^{1/2} \quad (6)$$

with M(r,t) the mass enclosed by the collapsing surface.

From Equations (5) and (6) we see that in order to satisfy the requirement of timelike world line completeness for a collapsing object, for which $U_s < 0$, it is necessary to dynamically enforce the 'no trapped surface condition', $\frac{2GM_s}{c^2R_s} < 1$. In the MECO model, this is accomplished by the non-gravitational force of outflowing radiation. At the comoving surface, the luminosity is $L_s = 4\pi R^2 q \quad > 0$, where

$$q = -\frac{(c^2dM/d\tau)_s}{4\pi R^2(\Gamma_s + U_s/c)} = \frac{L_\infty(1+z_s)^2}{4\pi R^2} \quad (7)$$

and the distantly observed luminosity is L_∞.

To guarantee the existence of sufficient internal radiation pressure, it is likely that a MECO must possess an equipartition magnetic dipole field. At the temperatures and compactness of stellar collapse, a pair plasma exists within such a field. In addition to the intrinsic resistance to collapse of magnetic flux (Thorne 1965), it has been shown (Pelletier & Markowith 1998) that the energy of magnetic perturbations in equipartition pair plasmas is preferentially expended in photon production rather than causing particle acceleration. Photon pressure varies $\propto B^4$, due to its dependence on pair density ($\propto B^2$) and synchrotron photon energy ($\propto B^2$). Lacking the pair plasma, the ratio of magnetic ($\propto B^2$) to gravitational stresses would be constant in a collapsing gas (e.g. Baumgarte & Shapiro 2003). With photon pressure capable of increasing more rapidly than gravitational stress, a secular equilibrium rate of collapse can be stabilized with the radiation temperature buffered near the pair production threshold. The stability of the rate of collapse is maintained by increased (decreased) photon pressure ($\propto B^4$) if the field is increased (decreased) by compression (expansion). An equipartition field also easily confines the pair plasma. Thus the collapse differs in a fundamental way from that of only weakly magnetic, radiation dominated polytropic gas or pressureless dust.

Strong recent evidence for equipartition magnetic fields in stellar collapse has been found for GRB021206 (Coburn & Boggs 2003) and strong residual fields much in excess of those expected from mere flux compression have been found in magnetars (Ibrahim, Swank & Parke 2003). Kluzniak and Ruderman (1998) have described the generation of $\sim 10^{17}$ G magnetic fields for nuclear densities via differential rotation in neutron stars. Other possibilities for producing extreme magnetic fields would include ferromagnetic phase transitions during the collapse (Haensel & Bonnazzola 1996) or the formation of quark condensates (Tatsumi 2000.)

Since distantly observed magnetic fields are reduced by $\sim 1+z$, a redshift of $z \sim 10^8$ would be needed for the MECO model with an equipartition field to accord with the magnetic moments we have found for GBHC, and also to account for AGN luminosity constraints (see the calculation for Sgr A* below). *Thus we are motivated by the SPOE and empirical observational constraints to look for solutions of the GR field equations that are consistent with objects in extremely redshifted, Eddington limited gravitational collapse.*

3. EDDINGTON LIMITED MECO

The two key proper time differential equations that control the behavior of the surface of an Eddington balanced, collapsing, radiating object are: (Hernandez Jr. & Misner 1966, Lindquist, Schwartz & Misner 1965, Misner 1965):

$$\frac{dU_s}{d\tau} = (\frac{\Gamma}{\rho + P/c^2})_s(-\frac{\partial P}{\partial R})_s - (\frac{GM}{R^2})_s \quad (8)$$

Where $M_s = (M + 4\pi R^3(P+q)/c^2)_s$ includes magnetic field energy in P and radiant energy in q and

$$\frac{dM_s}{d\tau} = -(4\pi R^2 Pc\frac{U}{c})_s - (L(\frac{U}{c} + \Gamma))_s \quad (9)$$

In Eddington limited steady collapse, the condition, $dU_s/d\tau \approx 0$, holds. With this condition, Equation (8), when integrated over the closed surface where the pressure is dominantly that of radiation, can be solved for the net outward flow of Eddington limit luminosity through the surface. Taking the escape cone factor of $27(GM_s/c^2R_s)^2/(1+z_s)^2$ into account, the outflowing (but not all escaping) surface luminosity, L, would be

$$L_{Edd}(outflow)_s = \frac{4\pi GM_s cR^2(1+z_{Edd,s})^3}{27\kappa R_g^2} \quad (10)$$

where $R_g = GM_s/c^2$ and κ is the plasma opacity. For simplicity, we have assumed here that the luminosity actually escapes from the MECO surface rather than after conveyance through a MECO pair photosphere. The end result is the same for distant observers. However the luminosity L_s that appears in Equations (8 - 9) is actually the net luminosity, which escapes through the photon sphere, and is given by $L_{Edd}(escape)_s = L_{Edd}(outflow)_s - L_{Edd}(fallback)_s = L_{Edd}(outflow)_s - L_{Edd}(outflow)_s(1 - 27R_g^2/(R(1+z_{Edd}))^2$ Thus in Equations (8) and (9), the L_s appearing there is given by

$$L_s = L_{Edd}(escape)_s = \frac{4\pi GM(\tau)_s c(1+z_{Edd,s})}{\kappa} \quad (11)$$

Due to the thermal buffering provided by the equipartition field and pair plasma we can examine a limiting

case for which MECO mean proper density varies slowly enough that the condition $U_s/c << 1/(1+z_s) \approx \Gamma_s$ also holds after a time, τ_{Edd}, that has elapsed in reaching the Eddington limited state. In this context from (9) we have that

$$\frac{c^2 dM_s}{d\tau} = -\frac{L_{Edd}(escape)_s}{1+z_s} = -\frac{4\pi GM(\tau)_s c}{\kappa} \qquad (12)$$

which can be integrated to give

$$M_s(\tau) = M_s(\tau_{Edd}) \exp\left((-4\pi G/\kappa c)(\tau - \tau_{Edd})\right) \qquad (13)$$

For example, for hydrogen opacity, $\kappa = 0.4$ cm^2/g, and $z = 10^8$, this yields a distantly observed MECO lifetime of $(1+z_s)\kappa c/4\pi G \sim 5 \times 10^{16}$ yr. The MECO state for GBHC is likely preceded by a much faster gravitational collapse of a stellar core. With a neutrino opacity some 10^{20} times smaller than that of photons, the lifetime of Eddington limited neutrino emissions would likely be minutes, at most. To stabilize the rate of collapse with magnetic pressure and synchrotron generated photons would require a photon luminosity reduced below the neutrino Eddington luminosity by the same factor of $\sim 10^{20}$, to a distantly observed $\sim 10^{32}$ erg/s. It is of interest to note that for this to correspond to a MECO object radiating at its local Eddington limit, a surface redshift of $z_s \sim 10^{7-8}$ would be required, which accords with our earlier arguments based on empirical magnetic field observational constraints.

4. THE QUIESCENT MECO

Distantly observed MECO luminosity is diminished by $1/(1+z)^2$ by gravitational redshift and by $27R_g^2/(R(1+z))^2 \sim 27/(4(1+z)^2)$ by a narrow escape cone at the photosphere of a pair atmosphere. Due to its negligible mass, we consider the pair atmosphere to be external to the MECO. It can be shown *ex post facto* to be radiation dominated such that $T^4/(1+z)$ is constant throughout. Estimates of luminosities, photosphere upper limit temperatures and photosphere redshifts can then be found from the Eddington balance requirements.

The fraction of luminosity from the MECO surface that escapes to infinity in Eddington balance is

$$(L_{Edd})_s = \frac{4\pi GM_s c(1+z)}{\kappa} = 1.27 \times 10^{38} m(1+z_s) \;\; erg/s \qquad (14)$$

where $m = M/M_\odot$. The distantly observed luminosity is:

$$L_\infty = \frac{(L_{Edd})_s}{(1+z_s)^2} = \frac{4\pi GM_s c}{\kappa(1+z_s)} = \frac{1.27 \times 10^{38} m}{(1+z_s)} \;\; erg/s \qquad (15)$$

By assuming that the escaping radiation is primarily thermal and that the photosphere temperature is T_p, the fraction that escapes to be distantly observed is:

$$L_\infty = \frac{4\pi R_g^2 \sigma T_p^4 27}{(1+z_p)^4} = 1.56 \times 10^7 m^2 T_p^4 \frac{27}{(1+z_p)^4} \;\; erg/s \qquad (16)$$

where $\sigma = 5.67 \times 10^{-5}$ erg/s/cm^2 and subscript p refers to conditions at the photosphere. Equations (15) and (16) yield:

$$T_\infty = T_p/(1+z_p) = \frac{2.3 \times 10^7}{(m(1+z_s))^{1/4}} \;\; K. \qquad (17)$$

To examine typical cases, a GBHC with $m = 10$ and $z \sim 10^8$ would have $T_\infty = 1.3 \times 10^5 K = 0.01$ keV, a luminosity, excluding spin-down contributions, of $L_\infty = 1.3 \times 10^{31} erg/s$, and a spectral peak at 220 A^0, in the photoelectrically absorbed deep UV. For an m=10^8 AGN, $T_\infty = 2300K$, and $L_\infty = 1.3 \times 10^{38} erg/s$ with a spectral peak in the near infrared at 1.2 micron. *(Sgr A*, with $m = 3 \times 10^6$, would have $T_\infty = 5500$ K and a 2.2 micron brightness of 6 mJy, just below the observational upper limit of 9 mJy (Reid et al. 2003).)* Since $T_\infty = T_p/(1+z_p)$, $T_p^4/(1+z_p) = T_s^4/(1+z_z)$ and $T_s \approx 6 \times 10^9$ K, we find that

$$T_p = T_s\left(\frac{T_s}{T_\infty(1+z_s)}\right)^{1/3} = 3.8 \times 10^{10} \frac{m^{1/12}}{(1+z_s)^{1/4}} \;\; K \qquad (18)$$

For a GBHC with $m = 10$ and $z_s = 10^8$, this yields a photosphere temperature of 4.6×10^8 K, from which $(1+z_p) = 3500$. An AGN with $m = 10^8$ would have a somewhat warmer photosphere at $T_p = 1.8 \times 10^9$ K, but with a red shift of 7.7×10^5.

Hence, although they are not black holes, passive MECO without accretion disks would (using any realistic opacity) have lifetimes much greater than a Hubble time and emit highly red shifted quiescent thermal spectra that may be quite difficult to observe.

5. THE HIGH STATE OF AN ACTIVELY ACCRETING MECO

From the viewpoint of a distant observer, accretion would deliver mass-energy to the MECO, which would then radiate most of it away. The contribution from the central MECO alone would be

$$L_\infty = \frac{4\pi GM_s c}{\kappa(1+z_s)} + \frac{\dot{m}_\infty c^2}{1+z_s}(e(1+z_s)-1) = 4\pi R_g^2 \sigma T_p^4 \frac{27}{(1+z_p)^4} \qquad (19)$$

where $e = E/m_0c^2 = 0.943$ is the specific energy per particle available after accretion disk flow to the marginally stable orbit radius, r_{ms}. Assuming that $\dot{m}_\infty$ is some fraction, f, of the Newtonian Eddington limit mass accretion rate, $4\pi GMc/\kappa$, then

$$1.27 \times 10^{38} \frac{m\eta}{1+z_s} = (27)(1.56 \times 10^7) m^2 \left(\frac{T_p}{1+z_p}\right)^4 \qquad (20)$$

where $\eta = 1 + f((1+z_s)e - 1)$ includes both quiescent and accretion contributions to the luminosity. Due to the extremely strong dependence on temperature of the density of pairs, it is likely that the photosphere temperature remains near the previously found $4.6 \times 10^8 K$. Then with $z_s = 10^8$, $m = 10$, and $f = 1$, we find $T_\infty = T_p/(1+z_p) = 1.3 \times 10^7$ K $= 1.1$ keV, and $(1+z_p) = 35$, which indicates considerable photospheric expansion. The MECO luminosity would be $L_\infty = 1.2 \times 10^{39}$ erg/s, which is approximately at the Newtonian Eddington limit. For comparison, the accretion disk outside the marginally stable orbit at r_{ms} (efficiency = 0.057) would produce only 6.8×10^{37} erg/s, with an inner disk temperature also 'ultrasoft' at ~ 1.1 keV.

Most photons escaping the photon sphere would depart with some azimuthal momentum on spiral trajectories that would eventually take them across and through the accretion disk. Thus a very large fraction of the soft photons would be subject to bulk comptonization in the

plunging region inside r_{ms}. This contrasts sharply with the situation for neutron stars where there is no comparably large plunging region. This accounts for the fact that hard x-ray spectral tails are comparatively much stronger for high state GBHC. Our preliminary calculations for photon trajectories randomly directed upon leaving the photon sphere indicate that this process would produce a power law component with photon index greater than 2.

6. DETECTING MECO

It may be possible to detect MECO in several ways. Firstly, for a red shift of $z \sim 10^8$, the quiescent luminosity of a GBHC MECO would be $\sim 10^{31} erg/s$ with $T_\infty \sim 0.01$ keV. This thermal peak in the strongly absorbed UV might be observable for very nearby or high galactic latitude GBHC, such as A0620-00 or XTE J1118+480. Secondly, At high state luminosities above $\sim 10^{36}$ erg/s, a central MECO would be a bright, small 'ultrasoft' central object that might be sharply eclipsed in deep dipping sources. Thirdly, a pair plasma atmosphere in an equipartition magnetic field should be virtually transparent to photon polarizations perpendicular to the magnetic field lines. The x-rays from the central MECO should exhibit some polarization that might be detectable. If GBHC MECO are the offspring of massive star supernovae, then they should be found all over the galaxy. Based upon our estimates of their quiescent temperatures, isolated GBHC MECO would be weak, polarized, EUV sources with a power-law tail in soft x-rays.

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